

CHIRAL SYMMETRY RESTORATION AS THE GEORGI VECTOR LIMIT*

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ABSTRACT

I discuss recent work done with Gerry Brown on chiral phase transition at high temperature and/or density described in terms of Georgi's vector limit. The notion of "mended symmetry" is suggested to play an important role in understanding the properties of hadrons in dense and/or hot matter before reaching the phase transition. It is shown that while the QCD vacuum in baryon-free space is resistant to "melting" up to the critical temperature, baryon-rich medium can induce the vacuum to become softer in temperature: the hadron masses drop faster in temperature when baryon matter is present.

Introduction

At low energies, various hadrons enter into certain physical processes without apparent relations between them. Thus for instance, the pion-nucleon scattering near threshold can be described with pions and nucleons alone or with pions, nucleons and the scalar meson σ with a large mass or pions, nucleons and the vector meson ρ . The common element is the spontaneously broken chiral symmetry which puts constraints so as to make all these diverse approaches give the same answer. At low energies, there is no obvious *raison-d'être* for the presence of the σ or the ρ . In this talk, I would like to argue that as one increases temperature and/or density of the hadronic system as in heavy-ion collisions or compact stars, the π , the σ , the ρ etc. come together into a pattern that reveals "hidden symmetries." The reasoning is very much like what Weinberg calls "mended symmetry"[1] and involves

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in particular the notion of the “vector limit” proposed by Georgi[2] some years ago. Both Weinberg and Georgi invoke the large N_c limit as the appropriate limit. In papers recently published by us[3] and to appear[4], we proposed that high T and/or high density would bring the matter to the corresponding limits without involving explicit large N_c assumption.

To illustrate the main point, I discuss three cases: 1) the role of the σ meson in nuclear physics; 2) the quark-number susceptibility at high temperature measured in lattice gauge calculations; 3) “cool” kaons observed in heavy-ion collisions.

Mended symmetry and the linear σ model

At zero temperature ($T = 0$) and zero density ($\rho = 0$), low-energy pion interactions are described by the nonlinear σ model. No low-mass scalars (*i.e.*, the σ) figure in the description and pions are derivatively coupled. There is of course a scalar associated with the trace anomaly of QCD

$$\chi \sim (G_{\mu\nu}^2)^{1/4} \quad (1)$$

where $G_{\mu\nu}$ is the gluon field tensor, but this involves a high excitation with its mass greater than the chiral scale $\Lambda_\chi \sim 1$ GeV and so decouples from the low-energy dynamics. Thus for instance, pion-nucleon scattering at low energy can be described by a Lagrangian that contains the nucleon and the pion field but no scalar particle.

However when one studies nucleon interactions in nuclei, one finds that there is a definite need for a scalar particle σ that gives the principal attraction between nucleons. More precisely this may be just a scalar field that interpolates 2π , 4π etc. correlations. This can be understood as follows.

Suppose that nuclear environments provide kinematic situation where higher energy (or momentum) scale is involved. This suggests imposing certain high energy constraints on the low-energy properties embodied in the current algebra Lagrangian of the nonlinear σ model. Weinberg has considered the constraints provided by Regge asymptotic properties and showed that certain symmetries are “mended” although the system is still in a broken phase. He thus identifies the quartet π , ϵ , ρ_0 and a_{10} (with the subscript 0 denoting the helicity-0 component) to satisfy a mended symmetry relation. Using this argument, Beane and van Kolck[5] suggested that such high-energy constraints bring the glueball field χ down in energy and turn it into a dilaton σ with its mass nearly degenerate with the pion mass[†]. The symmetry mended in this way – while the mass of the scalar is not quite degenerate with the pion – is then the $O(4)$ symmetry involving the quartet π^1, π^2, π^3 and σ . The resulting

[†]To be more precise, it is the quarkish component of the χ field that becomes a dilaton as clarified in [4]. The non-smooth (high-frequency) glueball component is integrated out and does not figure explicitly in low-energy dynamics.

Lagrangian is precisely the linear σ model, apart from a logarithmic potential term that arises from the incorporation of the QCD trace anomaly. This Lagrangian is then assumed to be applicable at some density not far from a critical point. We suggest that this model is applicable already at nuclear matter density with the scalar mass $m_\sigma \sim 400$ MeV, the meson that figures in Walecka's model and also in one-boson-exchange potentials.

I should mention that this work gives support to the scaling proposed by us[6] for the hadrons in medium.

Assuming that the Beane-van Kolck model describes a dense medium with its constants approaching the dilaton limit, an interesting question is how hadrons respond to temperature. We know from chiral perturbation theory (supported by lattice calculations) that the properties of hadrons in zero-density space do not change appreciably up to the chiral phase transition temperature $T_c \approx 140$ MeV. For instance, at low temperature, chiral perturbation theory (or the linear σ model) predicts that the pion decay constant has the temperature dependence

$$\frac{f_\pi(T)}{f_\pi} = 1 - \frac{T^2}{12f_\pi^2} + \dots \quad (2)$$

Now consider the Beane-van Kolck Lagrangian in the dilaton limit given by

$$\begin{aligned} L = & i\bar{Q} \not{\partial} Q + \frac{1}{2} \partial_\mu \vec{\pi} \partial^\mu \vec{\pi} + \frac{1}{2} \partial_\mu \sigma \partial^\mu \sigma - \frac{m^*}{f_\pi} \bar{Q} [\sigma - i\gamma_5 \vec{\pi} \cdot \vec{\tau}] Q \\ & + \frac{m_\sigma^{*2}}{16f_\pi^{*2}} (\sigma^2 + \vec{\pi}^2)^2 - \frac{m_\sigma^{*2}}{8f_\pi^{*2}} [(\sigma^2 + \vec{\pi}^2)^2 \ln(\sigma^2 + \vec{\pi}^2)/f_\pi^{*2}] \end{aligned} \quad (3)$$

where Q is the constituent quark field with mass m^* , σ the dilaton field with mass m_σ^* and the pion decay constant f_π^* is the in-medium constant given by the Brown-Rho scaling. With this Lagrangian, the temperature dependence of the pion decay constant is calculated to be[7]

$$\frac{f_\pi^*(T)}{f_\pi^*} = 1 - \frac{T^2}{4f_\pi^{*2}} + \dots \quad (4)$$

This shows that in the presence of baryonic matter, the pion decay constant drops substantially faster in temperature than in baryon-free space (2). (Note that at $\rho \approx \rho_0$, $f_\pi^*/f_\pi \approx 0.8$.) One can also show[7] that the BR scaling holds approximately in this model as a function of temperature. It would be interesting to check this property in lattice gauge calculations.

The Georgi vector limit

In a way analogous to the appearance of the low-mass scalar σ discussed above, vector mesons can also be brought into the low-energy sector. This involves the Georgi vector

symmetry and vector limit[2]. Consider two chiral flavors u(p) and d(own), with chiral symmetry $SU(2)_L \times SU(2)_R$. The standard way of looking at this symmetry is that it is realized either in Nambu-Goldstone (or Goldstone in short) mode, with $SU(2)_L \times SU(2)_R$ broken down spontaneously to $SU(2)_{L+R}$ or in Wigner-Weyl (or Wigner in short) mode with parity doubling. Georgi observes, however, that there is yet another way of realizing the symmetry which requires both Goldstone mode and Wigner mode to *co-exist*. Now the signature for *any* manifestation of the chiral symmetry is the pion decay constant f_π

$$\langle 0 | A_\mu^i | \pi^j(q) \rangle = i f_\pi q_\mu \delta^{ij} \quad (5)$$

where A_μ^i is the isovector axial current. The Goldstone mode is characterized by the presence of the triplet of Goldstone bosons, π^i with $i = 1, 2, 3$ with a non-zero pion decay constant. The Wigner mode is realized when the pion decay constant vanishes, associated with the absence of zero-mass bosons. In the latter case, the symmetry is realized in a parity-doubled mode. The Georgi vector symmetry we are interested in corresponds to the mode (5) co-existing with a triplet of scalars S^i with $f_S = f_\pi$ where

$$\langle 0 | V_\mu^i | S^j(q) \rangle = i f_S q_\mu \delta^{ij} \quad (6)$$

where V_μ^i is the isovector-vector current. In this case, the $SU(2) \times SU(2)$ symmetry is unbroken. At low T and/or low density, low-lying isovector-scalars are *not* visible and hence either the vector symmetry is broken in Nature with $f_S \neq f_\pi$ or they are “hidden” in the sense that they are eaten up by vector particles (à la Higgs). We are suggesting that as temperature and/or density rises to the critical value corresponding to the chiral phase transition, the symmetry characterized by

$$f_S = f_\pi \quad (7)$$

is restored with the isovector scalars making up the longitudinal components of the massive ρ mesons, which eventually get “liberated” at some high temperature (or density) from the vectors and become degenerate with the zero-mass pions at $T \gtrsim T_{\chi SR}$ where $T_{\chi SR} \sim T_c \sim 140$ MeV is the chiral transition temperature. The symmetry (7) with the scalars “hidden” in the massive vector mesons resembles Weinberg’s mended symmetry[1]. We shall reserve “Georgi vector limit” as the symmetry limit in which (7) holds together with $m_\pi = m_S = 0$.

The relevant Lagrangian to use for illustrating what we mean is the hidden gauge symmetric Lagrangian of Bando *et al* [8] which is valid below the chiral transition[‡],

$$\mathcal{L} = \frac{1}{2} f^2 \left\{ \text{Tr}(D^\mu \xi_L D_\mu \xi_L^\dagger) + (L \rightarrow R) \right\} + \kappa \cdot \frac{1}{4} f^2 \text{Tr}(\partial^\mu U \partial_\mu U^\dagger) + \dots \quad (8)$$

[‡]We will ignore the small quark masses and work in the chiral limit.

where $U = \xi_L \xi_R^\dagger$, $D^\mu \xi_{L,R} = \partial^\mu \xi_{L,R} - ig \xi_{L,R} \rho^\mu$, $\rho_\mu \equiv \frac{1}{2} \tau^a \rho_\mu^a$ and g stands for the hidden gauge coupling. The ellipsis stands for other matter fields and higher-derivative terms needed to make the theory more realistic. The ξ field can be parametrized as $\xi_{L,R} \equiv e^{iS(x)/f_S} e^{\pm i\pi(x)/f_\pi}$ with $S(x) = \frac{1}{2} \tau^a S^a(x)$ and $\pi(x) = \frac{1}{2} \tau^a \pi^a(x)$. The symmetry of the Lagrangian (8) is $(SU(2)_L \times SU(2)_R)_{global} \times G_{local}$ with $G \in SU(2)_V$. Setting $S(x) = 0$ corresponds to taking the unitary gauge in which case we are left with physical fields only (*i.e.*, no ghost fields).

At tree level, we get that

$$f_S = f, \quad f_\pi = \sqrt{1 + \kappa} f \quad (9)$$

and the $\rho\pi\pi$ coupling

$$g_{\rho\pi\pi} = \frac{1}{2(1 + \kappa)} g. \quad (10)$$

Going to the unitary gauge, one gets the KSRF mass relation

$$m_\rho = fg = \frac{1}{\sqrt{1 + \kappa}} f_\pi g = 2\sqrt{1 + \kappa} f_\pi g_{\rho\pi\pi}. \quad (11)$$

We know from experiments that at zero T (or low density), the κ takes the value $-\frac{1}{2}$ for which the KSRF relation is accurately satisfied. The symmetry (7) therefore is broken. The symmetry is recovered for $\kappa = 0$ in which case the second term of (8) that mixes L and R vanishes, thus restoring $SU(2) \times SU(2)$. In this limit, the hidden gauge symmetry swells to $G_L \times G_R$, and $\xi_{L,R}$ transform $\xi_L \rightarrow L \xi_L G_L^\dagger$, $\xi_R \rightarrow R \xi_R G_R^\dagger$. If the gauge coupling is not zero, then the ρ mesons are still massive and we have the mended symmetry (7). However if the gauge coupling vanishes, then the vector mesons become massless and their longitudinal components get liberated, giving the scalar massless multiplets $S(x)$. In this limit, the symmetry is the global $[SU(2)]^4$. Local symmetry is no longer present.

It is our proposal that in hot and/dense matter approaching the chiral restoration, the constant $\kappa \rightarrow 0$ and the gauge coupling $g \rightarrow 0$. For this purpose, we shall extrapolate a bit the results obtained by Harada and Yamawaki[9]. These authors studied the hidden gauge Lagrangian (8) to one loop order and obtained the β functions for the hidden gauge coupling g and the constant κ (in dimensional regularization)

$$\beta_g(g) = \mu \frac{dg}{d\mu} = -\frac{87 - a^2}{12} \frac{g^2}{(4\pi)^2}, \quad (12)$$

$$\beta_\kappa(a) = \mu \frac{da}{d\mu} = 3a(a^2 - 1) \frac{g^2}{(4\pi)^2} \quad (13)$$

with $a = \frac{1}{1 + \kappa}$. One notes that first of all, there is a nontrivial ultraviolet fixed point at $a = 1$ or $\kappa = 0$ and that the coupling constant g scales to zero as $\sim 1/\ln \mu$ in the ultraviolet limit. This perturbative result may not be realistic enough to be taken seriously – and certainly

cannot be pushed too high in energy-momentum scale but for the reason given below, we think it plausible that the Harada-Yamawaki results hold at least qualitatively as high T (or density) is reached. In fact we expect that the gauge coupling should fall off to zero much faster than logarithmically in order to explain what follows below.

Quark-number susceptibility

As a case for the role of the vector limit, we consider the lattice gauge calculations made by Gottlieb *et al* [11] of the quark-number susceptibility defined by

$$\chi_{\pm} = (\partial/\partial\mu_u \pm \partial/\partial\mu_d) (\rho_u \pm \rho_d) \quad (14)$$

where the $+$ and $-$ signs define the singlet (isospin zero) and triplet (isospin one) susceptibilities, μ_u and μ_d are the chemical potentials of the up and down quarks and

$$\rho_i = \text{Tr} N_i \exp \left[-\beta \left(H - \sum_{j=u,d} \mu_j N_j \right) \right] / V \equiv \langle \langle N_i \rangle \rangle / V \quad (15)$$

with N_i the quark number operator for flavor $i = u, d$. The χ_+ is in the ω -meson channel and the χ_- in the ρ -meson channel. Due to the $SU(2)$ symmetry, $\chi_+ = \chi_-$ as one can see in the lattice results. Since the singlet susceptibility is similar to the non-singlet one, we consider the latter.

One can classify the lattice results by roughly three temperature regimes. In the very low temperature regime, the χ_- is dominated by the ρ meson and is small. As the temperature moves toward the onset of the phase transition, the χ_- increases rapidly to near that of non-interacting quarks. This regime may be described in terms of constituent quarks. In RPA approximation of the constituent quark model as used by Kunihiro[12], the susceptibility below the critical temperature is

$$\chi = \frac{\chi_0}{1 + G_v \chi_0} \quad (16)$$

where G_v is the coupling of the constituent quark (denoted Q) to the vector meson ρ and χ_0 is the susceptibility for non-interacting quarks which at $T \approx T_{\chi SR}$ where the dynamical mass m_Q has dropped to zero has the value

$$\chi_0 \approx N_f T^2 \quad (17)$$

with N_f the number of flavors. In terms of the gauge coupling of (8), we have

$$G_v \approx \frac{g^2}{4m_\rho^2}. \quad (18)$$

As noted by Kunihiro in the NJL model, the rapid increase of the susceptibility can be understood by a steep drop in the vector coupling across the $T_{\chi SR}$. Let us see what we obtain with the hidden gauge symmetry Lagrangian (8). If we assume that the KSFR relation (11) holds at T near $T_{\chi SR} \approx 140$ MeV (the recent work by Harada *et al* [10] supports this assumption) and that $\chi_0 \approx 2T_{\chi SR}^2$ for $N_f = 2$, then we find

$$\chi(T_{\chi SR})/\chi_0(T_{\chi SR}) \approx \frac{1}{1 + \frac{1}{2}(\frac{T_{\chi SR}}{f_\pi})^2} \approx 0.47 \quad (19)$$

with $\kappa = 0$. Here we are assuming that f_π remains at its zero temperature value, 93 MeV, up to near $T_{\chi SR}$. The ratio (19) is in agreement with the lattice data at $T \lesssim T_{\chi SR}$.

Let us finally turn to the third regime, namely above $T_{\chi SR}$. It has been shown by Prakash and Zahed [13] that with increasing temperature, the susceptibility goes to its perturbative value which can be calculated with perturbative gluon-exchanges. The argument is made with the dimensional reduction at asymptotic temperatures, but it seems to apply even at a temperature slightly above $T_{\chi SR}$. We shall schematize the Prakash-Zahed argument using the dimensionally reduced model of Koch *et al* [14] which hints at the onset of the Georgi vector limit. In this model which exploits the “funny space” obtained by interchanging z and t , the helicity-zero state of the ρ meson is found to come out degenerate with the pion while the helicity \pm states are degenerate with each other. In finite temperature, z replaces T , so asymptotically in T , the configuration space with the new z becomes 2-dimensional with x and y . The ρ meson has gone massless and behaves like a (charged) photon with helicities ± 1 perpendicular to the plane. The helicity-zero state originating from the longitudinally polarized component of the ρ before it went massless now behaves as an isotriplet scalar. We identify this with the scalar $S(x)$ described above, a realization of the Georgi vector symmetry.

Let us assume then that the vector mesons have decoupled with $g = 0$. Going to the perturbative picture with quark-gluon exchanges, we take one-gluon-exchange potential of Koch *et al*,

$$V(r_t) = \frac{\pi}{m^2} \frac{4}{3} \bar{g}^2 T \sigma_{z,1} \sigma_{z,2} \delta(r_t) \quad (20)$$

with \bar{g} the color gauge coupling and $\delta(r_t)$ is the δ -function in the two-dimensional reduced space. Here $m = \pi T$ is the chiral mass of quark or antiquark as explained in [14]. Possible constant terms that can contribute to eq.(20) will be ignored as in [14]. In order to evaluate the expectation value of the $\delta(r_t)$, we note that the helicity-zero ρ -meson wave function in two dimensions is well approximated by $\psi_\rho \approx N e^{-r_t/a}$ with $a \approx \frac{2}{3}$ fm and the normalization $N^2 = \frac{2}{\pi a^2}$. For the helicity ± 1 ρ -mesons, $\sigma_{z,1} \sigma_{z,2} = 1$, so we find that the expectation value of V is

$$\langle V \rangle = \frac{8}{3} \frac{\bar{g}^2 T}{\pi^2 T^2 a^2}. \quad (21)$$

Now summing the ladder diagrams to all orders, we get

$$\frac{\chi}{\chi_0} = \left(1 + \frac{\langle V \rangle}{2\pi T}\right)^{-1}, \quad (22)$$

where the energy denominator $2\pi T$ corresponds to the mass of a pair of quarks.

The lattice calculations [11] use $6/\bar{g}^2 = 5.32$ which would give $\alpha_s = 0.07$ at scale of a^{-1} where a is the lattice spacing. Calculations use 4 time slices, so the renormalized \bar{g} is that appropriate to $a^{-1/4}$. Very roughly we take this into account by multiplying the above α_s by $\ln 4^2$; therefore using $\alpha_s \cong 0.19$. With this α_s and the above wave function, we find

$$\frac{\chi(T_{\chi SR}^+)}{\chi_0(T_{\chi SR}^+)} \approx 0.68. \quad (23)$$

This is just about the ratio obtained above $T_{\chi SR}$ in the lattice calculations. Remarkably the perturbative result (23) above T_c matches smoothly onto the HGS prediction (19) just below T_c . Neglecting logarithmic dependence of the gauge coupling constant, eq. (22) can be written as

$$\frac{\chi}{\chi_0}(T) \approx \frac{1}{1 + 0.46(T_c/T)^2} \quad (24)$$

which follows closely the lattice gauge results of Gottlieb *et al* [11]. We consider this an indication for the Georgi vector symmetry, with the induced flavor gauge symmetry in the hadronic sector ceding to the fundamental color gauge symmetry of QCD in the quark-gluon sector.

We should remark that to the extent that the screening mass obtained in [14] $m_\pi = m_S \approx 2\pi T$ is consistent with two non-interacting quarks and that the corresponding wave functions obtained therein are the same for the pion π and the scalar S , we naturally expect the relation (7) to hold.

Cool kaons in heavy-ion collisions

The vanishing of the hidden gauge coupling g can have a dramatic effect on the kaons produced in relativistic heavy-ion collisions. In particular, it is predicted that the kaons produced from quark-gluon plasma would have a component that has a temperature much lower than that for other hadrons. This scenario may provide an explanation of the recent preliminary data[15] on the 14.6 GeV collision (experiment E-814)

$$^{28}\text{Si} + \text{Pb} \rightarrow K^+(K^-) + X \quad (25)$$

which showed cool components with effective temperature of 12 MeV for K^+ and 10 MeV for K^- , which cannot be reproduced in the conventional scenarios employed in event generators. The latter give kaons of effective temperature ~ 150 MeV.

There are two points to keep in mind in understanding what is happening here. Firstly, the Brookhaven AGS experiments determined the freeze-out – the effective decoupling in the sense of energy exchange of pions and nucleons – at $T_{fo} \approx 140$ MeV[16]. This is essentially the same as the chiral transition temperature measured in lattice gauge calculations [17]. This suggests that the freeze-out for less strongly interacting particles other than the pion and the nucleon is at a temperature higher than $T_{\chi SR}$ and that the pion and nucleon freeze out at about $T_{\chi SR}$. This means that interactions in the interior of the fireball will be at temperature greater than $T_{\chi SR}$. At this temperature, the vector coupling g would have gone to zero, so the Georgi vector limit would be operative were it to be relevant. The second point is that the fireball must expand slowly. The slow expansion results because the pressure in the region for some distance above $T_{\chi SR}$ is very low [18], the energy in the system going into decondensing gluons rather than giving pressure. This results in an expansion velocity of $v/c \sim 0.1$. In the case of 15 GeV/N Si on Pb transitions, the fireball has been measured [19] through Hanbury-Brown-Twiss correlations of the pions to increase from a transverse size of $R_T(Si) = 2.5$ fm to $R_T = 6.7$ fm, nearly a factor of 3, before pions freeze out. With an expansion velocity of $v/c \sim 0.1$, this means an expansion time of $\sim 25 - 30$ fm/c. (The full expansion time cannot be measured from the pions which occur as a short flash at the end.)

In a recent paper, V. Koch[20] has shown that given a sizable effective attractive interaction between the K^+ and the nucleon at the freeze-out phase, a cool kaon component can be reproduced in the conditions specified above. He elaborated further on this in this meeting. We argue now that such an attractive interaction can result if the Georgi vector limit is realized.

The description by chiral perturbation theory[21, 22, 23, 24] of kaon nuclear interactions and kaon condensation in dense nuclear matter has shown that three mechanisms figure prominently in kaon-nuclear processes at low energy: (1) the ω meson exchange giving rise to repulsion for K^+N interactions and attraction for K^-N ; (2) the “sigma-term” attraction for both $K^\pm N$; (3) the repulsive “virtual pair term.” In effective chiral Lagrangians, the first takes the form, $\sim \pm \frac{1}{f^2} K^\dagger \partial_\mu K \bar{N} \gamma^\mu N$ for K^\pm , the second $\sim \frac{\Sigma_{KN}}{f^2} K^\dagger K \bar{N} N$ and the third term $\sim (\partial_\mu K)^2 \bar{N} N$. The vector-exchange gives the repulsive potential

$$V_{K+N} \cong \frac{1}{3} V_{NN} \cong 90 \text{ MeV } \frac{\rho}{\rho_0} \quad (26)$$

where ρ_0 is nuclear matter density. It is not explicit but this term is proportional to the hidden gauge coupling g^2 . As for the scalar attraction, it is mainly given by the “sigma term”

$$S_{K+N} \approx -\frac{\Sigma_{KN} \langle \bar{N} N \rangle}{2m_K f^2} \cong -45 \text{ MeV } \frac{\rho_s}{\rho_0} \quad (27)$$

where ρ_s is the scalar density and Σ_{KN} is the KN sigma term. This remains attractive for K^-N interactions. The virtual pair term removes[25], at zero temperature, about 60 % of the attraction (27). At low temperature, the net effect is therefore highly repulsive for K^+N interactions.

What happens as $T \rightarrow T_{\chi SR}$ is as follows. First of all, part of the virtual pair repulsion gets “boiled” off as discussed in [25]. More importantly, if the Georgi vector limit is relevant, then the vector mesons decouple with $g \rightarrow 0$, removing the repulsion (26). As a consequence, the residual attraction from the scalar exchange remains. The calculation of Koch[20] supports this scenario. Since the vector coupling is absent, both K^+ and K^- will have a similar cool component.

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